

How Magnetic is the Neutrino? *

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The existence of a neutrino magnetic moment implies contributions to the neutrino mass via radiative corrections. We derive model-independent “naturalness” upper bounds on the magnetic moments of Dirac and Majorana neutrinos, generated by physics above the electroweak scale. For Dirac neutrinos, the bound is several orders of magnitude more stringent than present experimental limits. However, for Majorana neutrinos the magnetic moment bounds are weaker than present experimental limits if μ_ν is generated by new physics at ~ 1 TeV, and surpass current experimental sensitivity only for new physics scales $> 10 - 100$ TeV. The discovery of a neutrino magnetic moment near present limits would thus signify that neutrinos are Majorana particles.

Keywords: Neutrino, magnetic moment, neutrino mass

1. Introduction

In the Standard Model (minimally extended to include non-zero neutrino mass) the neutrino magnetic moment is non-zero, but small, and is given by¹

$$\mu_\nu \approx 3 \times 10^{-19} \left(\frac{m_\nu}{1\text{eV}} \right) \mu_B, \quad (1)$$

where m_ν is the neutrino mass and μ_B is the Bohr magneton. An experimental observation of a magnetic moment larger than that given in Eq.(1) would thus be a clear indication of physics beyond the minimally extended Standard Model. Current laboratory limits are determined via neutrino-electron scattering at low energies, with $\mu_\nu < 1.5 \times 10^{-10} \mu_B$ ² and $\mu_\nu < 0.7 \times 10^{-10} \mu_B$ ³ obtained from solar and reactor experiments, respectively. A stronger limit can be obtained from constraints on energy loss from stars, $\mu_\nu < 3 \times 10^{-12} \mu_B$.⁴

It is possible to write down a simple relationship between the size of the neutrino mass and neutrino magnetic moment. If a magnetic moment is generated by physics beyond the Standard Model (SM) at an energy scale Λ , as in Fig. 1a, we can

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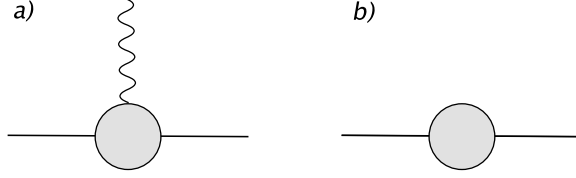


Fig. 1. a) Generic contribution to the neutrino magnetic moment induced by physics beyond the standard model. b) Corresponding contribution to the neutrino mass. The solid and wavy lines correspond to neutrinos and photons respectively, while the shaded circle denotes physics beyond the SM.

generically express its value as

$$\mu_\nu \sim \frac{eG}{\Lambda}, \quad (2)$$

where e is the electric charge and G contains a combination of coupling constants and loop factors. Removing the photon from the same diagram (Fig. 1b) gives a contribution to the neutrino mass of order

$$m_\nu \sim G\Lambda. \quad (3)$$

We thus have the relationship

$$m_\nu \sim \frac{\Lambda^2}{2m_e} \frac{\mu_\nu}{\mu_B} \sim \frac{\mu_\nu}{10^{-18} \mu_B} [\Lambda(\text{TeV})]^2 \text{ eV}, \quad (4)$$

which implies that it is difficult to simultaneously reconcile a small neutrino mass and a large magnetic moment.

However, it is well known that the naïve restriction given in Eq.(4) can be overcome via a careful choice for the new physics. For example, we may impose a symmetry to enforce $m_\nu = 0$ while allowing a non-zero value for μ_ν ,⁵⁻⁸ or employ a spin suppression mechanism to keep m_ν small.⁹ Note though, that these symmetries are typically broken by Standard Model interactions. By calculating contributions to m_ν generated by SM radiative corrections involving the magnetic moment interaction, we may thus obtain general, “naturalness” upper limits on the size of neutrino magnetic moments.

One possibility for allowing a large μ_ν while keeping m_ν small is due to Voloshin.⁵ The original version of this mechanism involved imposing an $SU(2)_\nu$ symmetry, under which the left-handed neutrino and antineutrino (ν and ν^c) transform as a doublet. The Dirac mass term transforms as a triplet under this symmetry and is thus forbidden, while the magnetic moment term is allowed as it transforms as a singlet. However, the $SU(2)_\nu$ symmetry is violated by SM gauge interactions. For Majorana neutrinos, the Voloshin mechanism may be implemented using flavor symmetries, such as those in Refs. 6–8. These flavor symmetries are not broken by SM gauge interactions but are instead violated by SM Yukawa interactions.^a

^aWe assume that the charged leptons masses are generated via the standard mechanism through

Below, we shall estimate the contribution to m_ν generated by SM radiative corrections involving the magnetic moment term. This allows us to set general, “naturalness” upper limits on the size of neutrino magnetic moments. For Dirac neutrinos, these limits are several orders of magnitude stronger than present experimental bounds.¹⁰ For Majorana neutrinos, however, the bounds are weaker.^{11,12}

2. Dirac Neutrinos

We assume that the magnetic moment is generated by physics beyond the SM at an energy scale Λ above the electroweak scale. In order to be completely model independent, the new physics will be left unspecified and we shall work exclusively with dimension $D \geq 4$ operators involving only SM fields, obtained by integrating out the physics above the scale Λ . We thus consider an effective theory that is valid below the scale Λ , respects the $SU(2)_L \times U(1)_Y$ symmetry of the SM, and contains only SM fields charged under these gauge groups.

We start by constructing the most general operators that could give rise to a magnetic moment operator, $\bar{\nu}_L \sigma^{\mu\nu} F_{\mu\nu} \nu_R$. Demanding invariance under the SM gauge group, we have the following 6D operators

$$\mathcal{O}_B^{(6)} = \frac{g'}{\Lambda^2} \bar{L} \tilde{\phi} \sigma^{\mu\nu} \nu_R B_{\mu\nu} , \quad \mathcal{O}_W^{(6)} = \frac{g}{\Lambda^2} \bar{L} \tau^a \tilde{\phi} \sigma^{\mu\nu} \nu_R W_{\mu\nu}^a . \quad (5)$$

where $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ and $W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a - g\epsilon_{abc} W_\mu^b W_\nu^c$ are the $U(1)_Y$ and $SU(2)_L$ field strength tensors, respectively, and g' and g are the corresponding couplings. The Higgs and left-handed lepton doublet fields are denoted ϕ and L , respectively, and $\tilde{\phi} = i\tau_2 \phi^*$.

After spontaneous symmetry breaking, both $\mathcal{O}_B^{(6)}$ and $\mathcal{O}_W^{(6)}$ contribute to the magnetic moment. Through loop diagrams these operators will generate contributions to the neutrino mass. For example, the diagram in Fig. 2 will generate a contribution to the neutrino mass operator, $\mathcal{O}_M^{(4)} = \bar{L} \tilde{\phi} \nu_R$. Using dimensional analysis, we estimate¹⁰

$$m_\nu \sim \frac{\alpha}{16\pi} \frac{\Lambda^2}{m_e \mu_B} \mu_\nu , \quad (6)$$

and thus

$$\mu_\nu \lesssim 3 \times 10^{-15} \mu_B \left(\frac{m_\nu}{1 \text{ eV}} \right) \left(\frac{1 \text{ TeV}}{\Lambda} \right)^2 . \quad (7)$$

If we take $\Lambda \simeq 1 \text{ TeV}$ and $m_\nu \lesssim 0.3 \text{ eV}$, we obtain the limit $\mu_\nu \lesssim 10^{-15} \mu_B$, which is several orders of magnitude stronger than current experimental constraints. Given the quadratic dependence upon Λ , this constraint becomes extremely stringent for Λ significantly above the electroweak scale.

Yukawa couplings to the SM Higgs boson. If the charged lepton masses are generated via a non-standard mechanism, SM Yukawa interactions do not necessarily violate flavor symmetries. However, such flavor symmetries must always be broken via some mechanism in order to obtain non-degenerate masses for the charged leptons.

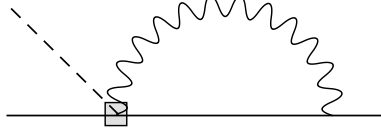


Fig. 2. Contribution to the 4D mass operator $\mathcal{O}_M^{(4)}$ due to insertions of the magnetic moment operators $\mathcal{O}_{B,W}^{(5)}$.

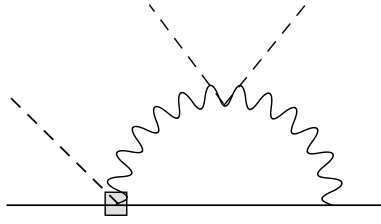


Fig. 3. Renormalization of the mass operator, $\mathcal{O}_M^{(6)}$, due to insertions of $\mathcal{O}_{B,W}^{(6)}$.

However, if Λ is not significantly larger than the EW scale, higher dimension operators are important, and their contribution to m_ν can be calculated in a model independent way. Through renormalization, both $\mathcal{O}_B^{(6)}$ and $\mathcal{O}_W^{(6)}$ will generate a contribution to the 6D neutrino mass operator

$$\mathcal{O}_M^{(6)} = \frac{1}{\Lambda^2} \bar{L} \tilde{\phi} \nu_R (\phi^\dagger \phi) , \quad (8)$$

via the diagrams in Fig. 3. Solving the renormalization group equations we find that for $\Lambda \gtrsim 1$ TeV,

$$\mu_\nu \lesssim 8 \times 10^{-15} \mu_B \left(\frac{m_\nu}{1 \text{ eV}} \right) , \quad (9)$$

in the absence of fine tuning.¹⁰

3. Majorana Neutrinos

We have seen above that the “naturalness” bounds on the magnetic moments of Dirac neutrinos are significantly stronger than present experimental limits. However, the analogous bounds for Majorana neutrinos are much weaker. The case of Majorana neutrinos is more subtle, due to the relative flavor symmetries of m_ν and μ_ν respectively. Majorana neutrinos cannot have diagonal magnetic moments, but are permitted non-zero transition moments. The transition magnetic moment $[\mu_\nu]_{\alpha\beta}$ is antisymmetric in the flavor indices $\{\alpha, \beta\}$, while the mass terms $[m_\nu]_{\alpha\beta}$ are symmetric. These different flavor symmetries play an important role in our limits, and are the origin of the difference between the magnetic moment constraints for Dirac and Majorana neutrinos.

As before, we write down the most general set of operators that can give rise to neutrino magnetic moment and mass terms, while respecting the SM gauge group. In the case of Majorana neutrinos, the lowest order contribution to the neutrino mass arises from the usual five dimensional operator containing Higgs and left-handed lepton doublet fields:

$$[O_M^{5D}]_{\alpha\beta} = (\overline{L}_\alpha^c \epsilon \phi) (\phi^T \epsilon L_\beta), \quad (10)$$

where $\epsilon = -i\tau_2$, $\overline{L}^c = L^T C$, C denotes charge conjugation, and α, β are flavor indices. The lowest order contribution to the neutrino magnetic moment arises from the following dimension seven operators,

$$[O_B]_{\alpha\beta} = g' (\overline{L}_\alpha^c \epsilon \phi) \sigma^{\mu\nu} (\phi^T \epsilon L_\beta) B_{\mu\nu}, \quad (11)$$

$$[O_W]_{\alpha\beta} = g (\overline{L}_\alpha^c \epsilon \phi) \sigma^{\mu\nu} (\phi^T \epsilon \tau^a L_\beta) W_{\mu\nu}^a, \quad (12)$$

and we also define a 7D mass operator as

$$[O_M^{7D}]_{\alpha\beta} = (\overline{L}_\alpha^c \epsilon \phi) (\phi^T \epsilon L_\beta) (\phi^\dagger \phi). \quad (13)$$

Operators O_M^{5D} and O_M^{7D} are flavor symmetric, while O_B is antisymmetric. The operator O_W is the most general 7D operator involving $W_{\mu\nu}^a$. However, as it is neither flavor symmetric nor antisymmetric it is useful to express it in terms of operators with explicit flavor symmetry, O_W^\pm , which we define as

$$[O_W^\pm]_{\alpha\beta} = \frac{1}{2} \{ [O_W]_{\alpha\beta} \pm [O_W]_{\beta\alpha} \}. \quad (14)$$

Our effective Lagrangian is therefore

$$\mathcal{L} = \frac{C_M^{5D}}{\Lambda} O_M^{5D} + \frac{C_M^{7D}}{\Lambda^3} O_M^{7D} + \frac{C_B}{\Lambda^3} O_B + \frac{C_W^+}{\Lambda^3} O_W^+ + \frac{C_W^-}{\Lambda^3} O_W^- + \dots \quad (15)$$

After spontaneous symmetry breaking, the flavor antisymmetric operators O_B and O_W^- generate a contribution to the magnetic moment interaction $\frac{1}{2} [\mu_\nu]_{\alpha\beta} \overline{\nu}_\alpha^c \sigma^{\mu\nu} \nu_\beta F_{\mu\nu}$, given by

$$\frac{[\mu_\nu]_{\alpha\beta}}{\mu_B} = \frac{2m_e v^2}{\Lambda^3} \left([C_B(M_W)]_{\alpha\beta} + [C_W^-(M_W)]_{\alpha\beta} \right), \quad (16)$$

where the Higgs vacuum expectation value is $\langle \phi^T \rangle = (0, v/\sqrt{2})$. Similarly, the operators O_M^{5D} and O_M^{7D} generate contributions to the Majorana neutrino mass terms, $\frac{1}{2} [m_\nu]_{\alpha\beta} \overline{\nu}_\alpha^c \nu_\beta$, given by

$$\frac{1}{2} [m_\nu]_{\alpha\beta} = \frac{v^2}{2\Lambda} [C_M^{5D}(M_W)] + \frac{v^4}{4\Lambda^3} [C_M^{7D}(M_W)]. \quad (17)$$

Below, we outline the radiative corrections to the neutrino mass operators (O_M^{5D} and O_M^{7D}) generated by the magnetic moment operators O_W^- and O_B . This allows us to determine constraints on the size of the magnetic moment in terms of the neutrino mass, using Eqs.(16) and (17). Our results are summarized in Table 1 below, where we have defined $R_{\alpha\beta} = m_\tau^2 / |m_\alpha^2 - m_\beta^2|$, with m_α being the masses of charged lepton masses. Numerically, one has $R_{\tau e} \simeq R_{\tau\mu} \simeq 1$ and $R_{\mu e} \simeq 283$.

Table 1. Summary of constraints on the magnitude of the magnetic moment of Majorana neutrinos. The upper two lines correspond to a magnetic moment generated by the O_W^- operator, while the lower two lines correspond to the O_B operator.

i) 1-loop, 7D	$\mu_{\alpha\beta}^W$	$\leq 1 \times 10^{-10} \mu_B \left(\frac{[m_\nu]_{\alpha\beta}}{1 \text{ eV}} \right) \ln^{-1} \frac{\Lambda^2}{M_W^2} R_{\alpha\beta}$
ii) 2-loop, 5D	$\mu_{\alpha\beta}^W$	$\leq 1 \times 10^{-9} \mu_B \left(\frac{[m_\nu]_{\alpha\beta}}{1 \text{ eV}} \right) \left(\frac{1 \text{ TeV}}{\Lambda} \right)^2 R_{\alpha\beta}$
iii) 2-loop, 7D	$\mu_{\alpha\beta}^B$	$\leq 1 \times 10^{-7} \mu_B \left(\frac{[m_\nu]_{\alpha\beta}}{1 \text{ eV}} \right) \ln^{-1} \frac{\Lambda^2}{M_W^2} R_{\alpha\beta}$
iv) 2-loop, 5D	$\mu_{\alpha\beta}^B$	$\leq 4 \times 10^{-9} \mu_B \left(\frac{[m_\nu]_{\alpha\beta}}{1 \text{ eV}} \right) \left(\frac{1 \text{ TeV}}{\Lambda} \right)^2 R_{\alpha\beta}$

3.1. $SU(2)$ Gauge Boson

3.1.1. 7D mass term — O_W

As the operator O_W^- is flavor antisymmetric, it must be multiplied by another flavor antisymmetric contribution in order to produce a flavor symmetric mass term. This can be accomplished through insertion of Yukawa couplings in the diagram shown in Fig. 4.¹¹ This diagram provides a logarithmically divergent contribution to the 7D mass term, given by¹¹

$$[C_M^{7D}(M_W)]_{\alpha\beta} \simeq \frac{3g^2}{16\pi^2} \frac{m_\alpha^2 - m_\beta^2}{v^2} \ln \frac{\Lambda^2}{M_W^2} [C_W^-(\Lambda)]_{\alpha\beta}, \quad (18)$$

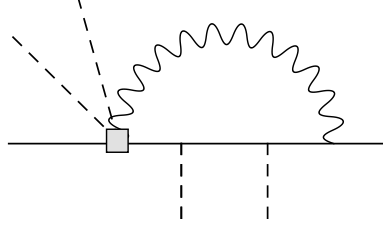
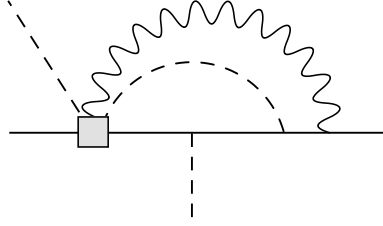
where m_α are the charged lepton masses, and the exact coefficient has been computed using dimensional regularization. Using this result, together with Eqs. (16) and (17), leads to bound (i) in Table 1.

3.1.2. 5D mass term — O_W

The neutrino magnetic moment operator O_W^- will also contribute to the 5D mass operator via two-loop diagrams, as shown in Fig. 5.¹² As with the diagrams in Fig. 4, we require two Yukawa insertions in order to obtain a flavor symmetric result. Using dimensional analysis, we estimate¹²

$$[C_M^{5D}(\Lambda)]_{\alpha\beta} \simeq \frac{g^2}{(16\pi^2)^2} \frac{m_\alpha^2 - m_\beta^2}{v^2} [C_W^-(\Lambda)]_{\alpha\beta}. \quad (19)$$

This leads to bound (ii) in Table 1. Compared to 1-loop (7D) case of Eq. (18), the 2-loop (5D) mass contribution is suppressed by a factor of $1/16\pi^2$ arising from the additional loop, but enhanced by a factor of Λ^2/v^2 arising from the lower operator dimension. Thus, as we increase the new physics scale, Λ , this two-loop constraint rapidly becomes more restrictive. The “crossover” scale between the two effects occurs at $\sim 10 \text{ TeV}$.

Fig. 4. Contribution of O_W^- to the 7D neutrino mass operator.Fig. 5. Representative contribution of O_W^- to the 5D neutrino mass operator.

3.2. Hypercharge Gauge Boson

3.2.1. 7D mass term — O_B

If we insert O_B in the diagram in Fig. 4, the contribution vanishes, due to the $SU(2)$ structure of the graph. Therefore, to obtain a non-zero contribution to O_M^{7D} from O_B we require the presence of some non-trivial $SU(2)$ structure. This can arise, for instance, from a virtual W boson loop as in Fig. 6.¹¹ This mechanism gives the leading contribution of the operator O_B to the 7D mass term. The O_B and O_W contributions to the 7D mass term are thus related by

$$\frac{(\delta m_\nu)^B}{(\delta m_\nu)^W} \approx \frac{\alpha}{4\pi} \frac{1}{\cos^2 \theta_W}, \quad (20)$$

where θ_W is the weak mixing angle and where the factor on the RHS is due to the additional $SU(2)_L$ boson loop. This additional loop suppression for the O_B contribution results in a significantly weaker neutrino magnetic moment constraint than that obtained above for O_W^- . The corresponding limit is shown as bound (iii) in Table 1.

3.2.2. 5D mass term — O_B

However, the leading contribution of O_B to the 5D mass term arises from the same 2-loop diagrams (Fig. 5) that we discussed in connection with the O_W^- operator. Therefore, the contribution to the 5D mass term is the same as that for O_W , except

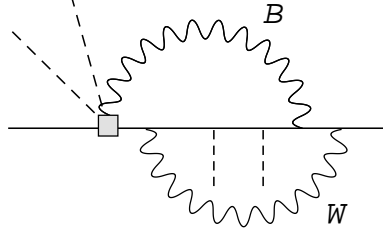


Fig. 6. Representative contribution of O_B to the 7D neutrino mass operator at two loop order.

for a factor of $(g'/g)^2 = \tan^2 \theta_W$. We thus obtain¹²

$$[C_M^{5D}(\Lambda)]_{\alpha\beta} \simeq \frac{g'^2}{(16\pi^2)^2} \frac{m_\alpha^2 - m_\beta^2}{v^2} [C_B(\Lambda)]_{\alpha\beta} \quad , \quad (21)$$

corresponding to bound (iv) in Table. 1. Importantly, this is the strongest constraint on the O_B contribution to the neutrino magnetic moment for any value of Λ , and the most general of our bounds on $\mu_\nu^{\text{Majorana}}$.¹²

3.3. Comparison with experimental limits

The best laboratory limit on μ_ν , obtained from scattering of low energy reactor neutrinos is, “ μ_e ” $< 0.7 \times 10^{-10} \mu_B$.³ Note that this limit applies to both $\mu_{\tau e}$ and $\mu_{\mu e}$, as the flavor of the scattered neutrino is not detected in the experiment. Taking the neutrino mass to be $m_\nu \lesssim 0.3$ eV (as implied by cosmological observations, e.g. Ref. 13), bound (iv) in Table. 1 gives

$$\begin{aligned} \mu_{\tau\mu}, \mu_{\tau e} &\lesssim 10^{-9} [\Lambda(\text{TeV})]^{-2} \\ \mu_{\mu e} &\lesssim 3 \times 10^{-7} [\Lambda(\text{TeV})]^{-2} . \end{aligned} \quad (22)$$

For Majorana neutrinos we thus conclude that if $\mu_{\mu e}$ is dominant over the other flavor elements, an experimental discovery near the present limits (e.g., at $\mu \sim 10^{-11} \mu_B$) would imply $\Lambda \lesssim 100$ TeV. However, this would become $\Lambda \lesssim 10$ TeV in any model in which all element of $\mu_{\alpha\beta}$ have similar size.

4. Conclusions

We have discussed radiative corrections to the neutrino mass arising from a neutrino magnetic moment coupling. Expressing the magnetic moment in terms of effective operators in a model independent fashion required constructing operators containing the $SU(2)_L$ and hypercharge gauge bosons, O_W and O_B respectively, rather than working directly with the electromagnetic gauge boson. We then calculated μ_ν naturalness bounds arising from the leading order contributions to neutrino mass term, for both Dirac and Majorana neutrinos. For Dirac neutrinos we found

$$\mu_\nu^{\text{Dirac}} \lesssim 3 \times 10^{-15} \mu_B \left(\frac{m_\nu}{1 \text{ eV}} \right) \left(\frac{1 \text{ TeV}}{\Lambda} \right)^2 , \quad (23)$$

while the most general naturalness bound on the size of the Majorana neutrino magnetic moment is

$$\mu_{\alpha\beta}^{\text{Majorana}} \leq 4 \times 10^{-9} \mu_B \left(\frac{[m_\nu]_{\alpha\beta}}{1 \text{ eV}} \right) \left(\frac{1 \text{ TeV}}{\Lambda} \right)^2 \left| \frac{m_\tau^2}{m_\alpha^2 - m_\beta^2} \right|. \quad (24)$$

These limits can only be evaded in the presence of fine tuning.

The limit on the the magnetic moments of Dirac neutrinos is thus considerably more stringent than for Majorana neutrinos. This is due to the different flavor symmetries involved, since in the Majorana case we require the insertion of Yukawa couplings to convert a flavor antisymmetric (magnetic moment) operator into a flavor symmetric (mass) operator. Our results implies that an experimental discovery of a magnetic moment near the present limits would signify (i) neutrinos are Majorana fermions and (ii) new lepton number violating physics responsible for the generation of μ_ν arises at a scale Λ which is well below the see-saw scale.

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